

The X-ray Pulsar M82 X-2 on its Propeller Line

D. M. Christodoulou,^{1,2★} D. Kazanas,^{3†} and S. G. T. Laycock^{1,4‡}

¹Lowell Center for Space Science and Technology, University of Massachusetts Lowell, Lowell, MA, 01854, USA

²Department of Mathematical Sciences, University of Massachusetts Lowell, Lowell, MA, 01854, USA

³NASA Goddard Space Flight Center, Laboratory for High-Energy Astrophysics, Code 663, Greenbelt, MD 20771, USA

⁴Department of Physics & Applied Physics, University of Massachusetts Lowell, Lowell, MA, 01854, USA

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ABSTRACT

NuSTAR has detected pulsations from the ultraluminous X-ray source X-2 in M82 and archival *Chandra* observations have given us a good idea of its duty cycle. The newly discovered pulsar exhibited at least 4 super-Eddington outbursts in the past 15 years but, in its lowest-power state, it radiates just below the Eddington limit and its properties appear to be typical of high-mass X-ray binaries. M82 X-2 has been described as a common neutron star with a 1 TG magnetic field that occasionally accretes above the Eddington rate and as a magnetar-like pulsar with a 10-100 TG magnetic field that reaches above the quantum limit. We argue in favor of the former interpretation. Using standard accretion theory and the available observations, we calculate the stellar magnetic field of this pulsar in two independent ways and we show that it cannot exceed 3 TG in either case. We discuss the implications of our results for other ultraluminous X-ray sources that occasionally exhibit similar powerful outbursts.

Key words: accretion, accretion disks—stars: magnetic fields—stars: neutron—X-rays: binaries—X-rays: individual (NuSTAR J095551+6940.8)

1 INTRODUCTION

Ultraluminous X-ray (ULX) sources are extragalactic compact accreting objects characterized by super-Eddington luminosities ($L_X \sim 10^{39-41}$ erg s^{−1}) and unusual soft X-ray spectra with blackbody emission around $\lesssim 0.3$ keV and a downturn above ~ 5 keV (Gladstone et al. 2009; Feng & Soria 2011; Motch et al. 2014). The extreme luminosities can be understood either as emission at the Eddington limit ($L_X \lesssim L_{\text{Edd}} = 1.3 \times 10^{38}$ M/M_⊙ erg s^{−1}) from intermediate-mass ($M \sim 10^{2-4}$ M_⊙) black holes or as anisotropic emission with $L_X > L_{\text{Edd}}$ from stellar-mass black holes and neutron stars (Soria 2007; Medvedev & Poutanen 2013; Bachetti et al. 2014; Motch et al. 2014; Pasham et al. 2014).

The latter interpretation is gaining favor from recent observations:

(a) Luangtip et al. (2014) analyzed a *Chandra* sample of nearby ULX sources and found a change in the spectral index around $L_X \sim 2 \times 10^{39}$ erg s^{−1} that may indicate a tran-

sition to the super-Eddington accretion regime by 10 M_⊙ black holes.

(b) Motch et al. (2014) determined that the ULX source P13 in NGC7793 harbors a stellar-mass black hole with $M < 15$ M_⊙.

(c) Bachetti et al. (2014) determined that the ULX source X-2 in M82 harbors a pulsar with spin period $P_S = 1.3725$ s and spinup rate $\dot{P}_S = -2 \times 10^{-10}$ s s^{−1}.

In all of these cases, the observations indicate that the X-ray sources radiate above their Eddington limits, a result that makes uncomfortable many researchers.

The magnetic field of NGC7793 P13 has not been estimated, but in the case of M82 X-2, Bachetti et al. (2014) used their measurement of the accretion torque to obtain a modest value of the magnetic field $B \approx 1$ TG. In contrast, Ekşi et al. (2015) argued that the accretion torque is very small near spin equilibrium and that leads to an underestimate of the magnetic field which in M82 X-2 must be of magnetar strength ($B \sim 100$ TG) (see also Ho et al. 2014; Klus et al. 2014; Lyutikov 2014; Dall’Osso et al. 2015; Tsygankov et al. 2016; Dall’Osso et al. 2016). We, on the other hand, find that the magnetic field implied from the *Chandra* measurements of M82 X-2 (Brightman et al. 2016; Tsygankov et al. 2016) is modest ($B \approx 3$ TG), a result that is in agreement with the independent determination

★ E-mail: dimitris_christodoulou@uml.edu

† E-mail: demos.kazanas@nasa.gov

‡ E-mail: silas_laycock@uml.edu

of [Brightman et al. \(2016\)](#) based on *NuSTAR* observations (a high-energy cutoff at 14_{-3}^{+5} keV). No unusual physics is necessary to obtain this estimate, the result is derived in two different ways from standard accretion theory and the measured values of P_S , \dot{P}_S , and minimum L_X in the lowest-power states. More importantly, such a modest B -value is consistent with isotropic X-ray emission near the Eddington limit when the pulsar lies in its low-power accreting state (on the so-called “propeller line;” [Christodoulou et al. 2016](#)). This consistency implies that M82 X-2 is a common pulsar whose propeller line lies near the Eddington limit ($L_X \approx L_{\text{Edd}}$) and that occasionally bursts anisotropically in our direction, an explanation also favored by [Bachetti et al. \(2014\)](#). On the other hand, [Tsygankov et al. \(2016\)](#) and [Dall’Osso et al. \(2016\)](#) have interpreted the low- L_X states of M82 X-2 as quiescent magnetospheric emission that occurs when the system straddles the [Corbet \(1996\)](#) gap. We show below that this interpretation is not supported by the observations.

The above results are already present in the equations of accretion on to compact objects ([Elsner & Lamb 1977](#); [Ghosh & Lamb 1979](#); [Wang 1987, 1995, 1996](#); [Frank et al. 2002](#)), but the physics is obscured by the multiple scalings and some minor inconsistencies in the definitions of the physical quantities involved. We thus derive and discuss the accreting pulsar’s minimum isotropic luminosity and surface magnetic field in § 2 without scaling the various physical variables, and we discuss our conclusions for M82 X-2 and other ULX sources in § 3 below.

2 EQUATIONS FOR ACCRETION ON TO A COMPACT OBJECT

We adopt Gaussian units for electromagnetic quantities ([Jackson 1962](#)) and we consider accretion from a disk-like structure that is formed around a compact object. As in [Pringle & Rees \(1972\)](#), we assume that on the propeller line the inner edge of the disk reaches inward to corotation (eqs. [1] and [8] below) and this condition is sufficient to determine the lowest emitted X-ray luminosity due to accretion. But we also consider the location of the magnetospheric radius (eqs. [2] and [6] below) in order to additionally derive an estimate of the stellar magnetic field as was done by [Ghosh & Lamb \(1979\)](#) and [Frank et al. \(2002\)](#).

We introduce below three control parameters in order to bridge the differences between many calculations carried out in the past: (a) parameter ξ specifies the location of the disk’s inner edge relative to the location of the spherical magnetospheric radius ([Ghosh & Lamb 1979](#); [Wang 1996](#)); (b) parameter α unifies two different definitions of the magnetic moment μ (e.g., [Wang 1996](#); [Ekşi et al. 2015](#)); and (c) η is the efficiency of converting accretion power to X-ray luminosity.

2.1 Accretion Torque

The accretion torque at the inner edge of the disk can be written as

$$\tau_d = \dot{M} \sqrt{GMR_d} \quad (1)$$

$$= \dot{M} \sqrt{GM(\xi r_m)}, \quad (2)$$

where \dot{M} is the mass accretion rate at the inner edge of the disk R_d , G is the gravitational constant, M is the mass of the compact object, and r_m is the spherical magnetospheric (Alfvén) radius. We assume that the accreted matter is forced to follow magnetic field lines at the cylindrical radius R_d and we adopt the relation $R_d \equiv \xi r_m$ ([Wang 1996](#)).

The accretion torque can be measured ([Bachetti et al. 2014](#)) as it can also be expressed in terms of the transfer of angular momentum onto the compact object during spinup:

$$|\tau_*| = |\dot{\mathcal{L}}_*| = \frac{2\pi I_*}{P_S^2} |\dot{P}_S|, \quad (3)$$

where $\mathcal{L}_* = 2\pi I_* / P_S$ and I_* are the angular momentum and the moment of inertia, respectively, of the compact object with spin period P_S and derivative \dot{P}_S . Once a nonzero τ_* has been measured, the torque in eq. (1) can no longer be adjusted: any attempt to lower τ_d using the “fastness parameter” ([Ekşi et al. 2015](#); [Dall’Osso et al. 2015](#)) leads to a corresponding increase in \dot{M} that maintains conservation of angular momentum. This pushes L_X higher and the increased luminosity no longer describes the propeller line.

2.2 Magnetospheric Radius

Calculations involving the spherical magnetospheric (Alfvén) radius r_m in the literature ([Elsner & Lamb 1977](#); [Ghosh & Lamb 1979](#); [Wang 1996](#); [Frank et al. 2002](#); [Ekşi et al. 2015](#)) suffer from inconsistencies in the definitions of the magnetic moment μ and the ram pressure \mathcal{P}_{ram} of the inflowing matter. Eq. (11) of [Ghosh & Lamb \(1979\)](#) agrees with eq. (6) of [Elsner & Lamb \(1977\)](#) only if $\mathcal{P}_{ram} = \rho v^2/2$ in the former calculation, where ρ is the mass density and v is the inflow speed of matter at spherical radius r . In both cases, the magnetic moment is defined as $\mu = Br^3$ on dimensional grounds (see also [Wang 1996](#)), where B is the magnetic field at r ; but other authors use $\mu = Br^3/2$ instead (e.g., [Ekşi et al. 2015](#)). The factor of 1/2 alters the definition of r_m but the modification has not been implemented by other authors who continue to use the classical formula for $r_m(\mu)$. Thus, various factors of 1/2 propagate in the calculations but they cause small differences in the coefficients of the final results because of the steep dependence of $B^2(r)$ on $1/r^6$. For example, using the 1/2 in the definition of μ would expand $r_m(\mu)$ by a factor of $16^{1/7} = 1.486$. We account for this discrepancy below by introducing the control parameter α in the definition of μ .

We adopt the following definitions: for the ram pressure, $\mathcal{P}_{ram} = \rho v^2/2$ (consistent with Bernoulli’s equation; [Wang 1996](#); [Menou et al. 1999](#)); and for the magnetic moment, $\mu = \alpha Br^3$, where α can be taken as 1 or 1/2. Then $\mu_*/\alpha = Br^3 = B_* R_*^3$, where asterisks indicate the quantities at the surface of the compact object, and the magnetic pressure at radius r can be written as $\mathcal{P}_{mag} = B^2/(8\pi) = \mu_*^2/(8\pi\alpha^2 r^6)$. As usual (e.g., [Frank et al. 2002](#)), $\rho|v| = \dot{M}/(4\pi r^2)$ and $|v|$ is close to the free-fall speed $v_{ff} = \sqrt{2GM/r}$, in which case the balance between the two pressures at $r = r_m$ determines the spherical Alfvén radius:

$$r_m = \left(\frac{1}{\alpha^4} \frac{\mu_*^4}{2GM\dot{M}^2} \right)^{1/7}. \quad (4)$$

The classical formula is recovered for $\alpha = 1$. It is convenient

to replace \dot{M} in terms of the X-ray luminosity L_X , which is assumed to be a fraction η of the total accretion power, i.e.,

$$L_X = \eta \left(\frac{GM\dot{M}}{R_*} \right), \quad (5)$$

in which case eq. (4) takes the form

$$r_m = \left(\frac{\eta^2}{2\alpha^4} \frac{GM\mu_*^4}{L_X^2 R_*^2} \right)^{1/7}. \quad (6)$$

2.3 Minimum Isotropic X-Ray Luminosity

Assuming that $|\tau_*| = \tau_d$, we get from eqs (1), (3), and (5):

$$L_X = \eta \left(\frac{2\pi I_*}{R_* P_S^2} |\dot{P}_S| \right) \sqrt{\frac{GM}{R_d}}. \quad (7)$$

If we ask for the inflowing matter that reaches R_d to also corotate with the magnetic field lines, then the inner-disk radius R_d must be equal to the corotation radius r_c :

$$R_d = r_c \equiv \left(\frac{GM P_S^2}{4\pi^2} \right)^{1/3}. \quad (8)$$

Eliminating R_d between eqs. (7) and (8), we find that the minimum X-ray luminosity on the propeller line is

$$L_X = \eta \left(2\pi I_* |\dot{P}_S| \right) \left(\frac{2\pi}{P_S^7} \frac{GM}{R_*^3} \right)^{1/3}. \quad (9)$$

The same equation (without the η) was also derived by Galache et al. (2008) who used it to obtain estimates of the minimum X-ray luminosity in a large sample of Magellanic high-mass X-ray binaries with measured P_S and \dot{P}_S values.

It is convenient to rewrite the result in terms of the angular velocity of the compact object $\Omega_S \equiv 2\pi/P_S$ and the Keplerian angular velocity on the surface $\Omega_{K*} \equiv \sqrt{GM/R_*^3}$:

$$L_X = \eta I_* (\Omega_S \Omega_{K*}^2)^{1/3} |\dot{\Omega}_S|. \quad (10)$$

The cubic root in this equation is the geometric mean of Ω_S and Ω_{K*}^2 . Thus, the luminosity $L_X \propto |\dot{\Omega}_S|$ and the proportionality constant is weighted twice by Ω_{K*} , the larger of the two frequencies in the geometric mean, which also expresses the depth of the gravitational potential well of the compact object. Despite that, it is well-known that the geometric mean favors the smaller values, therefore it is the spin frequency Ω_S that dominates the proportionality constant. This implies that a source that is spun up by accretion will increase its power output only marginally ($L_X \propto \Omega_S^{1/3}$) and corotation will move inward very slowly. For a compact object spinning near breakup ($\Omega_S = \Omega_{K*}$), corotation is near the surface of the object ($r_c \approx R_*$) and eq. (10) reduces to $L_X = \eta I_* \Omega_{K*} |\dot{\Omega}_K|$ which is simply the rate of radiative loss of a fraction η of the rotational kinetic energy $E_{rot} = I_* \Omega_{K*}^2 / 2$ that reaches the surface.

2.4 Stellar Magnetic Field

We set again $|\tau_*| = \tau_d$ and we combine eqs (2), (3), (5), (6), and (9) to get a relation that involves the stellar magnetic moment. We find that

$$\mu_* = \alpha (2\pi^2 \xi^7)^{-1/4} \sqrt{GMI_* |\dot{P}_S|}, \quad (11)$$

or, using $\mu_* = \alpha B_* R_*^3$, that the surface magnetic field is independent of α :

$$B_* = (2\pi^2 \xi^7)^{-1/4} \sqrt{\frac{GMI_*}{R_*^6}} |\dot{P}_S|. \quad (12)$$

The results in eqs. (11) and (12) do not depend on the precise amount of the inflowing accretion power that is converted to X-rays. Therefore, the efficiency of the conversion process (expressed by η in eq. [5]) does not enter in the calculation of B_* . This result makes physical sense. Eq. (12) also makes sense: it relates the stellar magnetic field B_* to nonmagnetic physical variables (see below), therefore B_* is independent of α .

Written in terms of Ω_S and Ω_{K*} , eq. (12) becomes

$$B_* = \left(\frac{\xi^7}{2} \right)^{-1/4} \left(\frac{\Omega_{K*}}{\Omega_S} \right) \sqrt{\frac{I_*}{R_*^3} |\dot{\Omega}_S|}. \quad (13)$$

This equation relates the surface magnetic pressure to the $R\phi$ component of the stress tensor $|\rho v_R v_K|_{R_d}$ at R_d , where $\rho(R_d) |v_R(R_d)| = \dot{M} / (4\pi R_d^2)$ and $v_K(R_d) = (GM/R_d)^{1/2}$: squaring both sides of eq. (13), we find that

$$\frac{B_*^2}{8\pi} = (2\xi^7)^{-1/2} \left(\frac{\Omega_{K*}}{\Omega_S} \right)^2 \left(\frac{R_d}{R_*} \right)^3 |\rho v_R v_K|_{R_d}. \quad (14)$$

This is because the torque balance $|\tau_*| = \tau_d$ implies that

$$\frac{I_*}{R_*^3} |\dot{\Omega}_S| = 4\pi |\rho v_R v_K|_{R_d} \left(\frac{R_d}{R_*} \right)^3. \quad (15)$$

Eq. (14) shows that the stress tensor at R_d is strongly rescaled by both the spin frequency of the compact object (ratio $(\Omega_{K*}/\Omega_S)^2 \gg 1$) and by the scale of the accretion flow (ratio $(R_d/R_*)^3 \gg 1$) in order to match the magnetic pressure on the surface of the compact object. Furthermore, eq. (14) reveals the following two notable accretion properties: (a) If the corotation radius R_d is smaller, then the magnetic field is weaker and vice versa, precisely as expected. (b) If the compact object is spun up (Ω_S increases), then the magnetic field appears to be weaker and vice versa, again as expected. The two effects do not work against one another because the direction in which the two ratios move during spinup or spindown is the same.

Let us now imagine that the compact object is spun up from accretion at a certain moment in time. Then P_S decreases, Ω_S increases, and corotation pushes inward (eq. [8]). This results in a decrease of both ratios in eq. (14) and then, since B_* remains unchanged on the surface, the $R\phi$ component of the stress tensor at R_d in eq. (14) must necessarily increase. Since the condition $|\tau_*| = \tau_d$ and eqs. (1) and (3) imply that $\dot{M} \propto 1/(P_S^2 R_d^{1/2})$, then \dot{M} also increases and this leads to an increase in the X-ray luminosity of the propeller line.

2.5 Another Form of the X-ray Luminosity on the Propeller Line

Eliminating $I_* |\dot{P}_S|$ between eqs. (9) and (11), we obtain the minimum luminosity on the propeller line (where $r_c = \xi r_m$ as in eq. (5) of Stella et al. 1986), that is

$$L_X = \left(\frac{\eta}{\alpha^2} \right) \left(\frac{\xi^7}{2} \right)^{1/2} \left(\frac{\Omega_S}{\Omega_{K*}} \right)^{4/3} \frac{\mu_*^2 \Omega_S}{R_*^3}. \quad (16)$$

The result of [Stella et al. \(1986\)](#) is reproduced for $\eta = 1 = \alpha$, $\xi = 0.5$ (which is reasonable for $L_X \ll L_{\text{Edd}}$), and canonical pulsar parameters. For these choices, the leading coefficient in eq. (16) is $1.4 \times 10^{37} \text{ erg s}^{-1}$, only by a factor of $1/\sqrt{2}$ lower than $2 \times 10^{37} \text{ erg s}^{-1}$ in [Stella et al. \(1986\)](#).

The propeller line, eq. (16) with $\xi = 0.5$, can be written in the physical form

$$L_X = \left(\frac{\eta}{16}\right) B_*^2 R_*^2 v_* \left(\frac{\Omega_S}{\Omega_{K*}}\right)^{4/3}, \quad (17)$$

where $v_* \equiv \Omega_S R_* = 2\pi R_*/P_S$. The interpretation of this equation is straightforward: The isotropic power output that emerges from the surface of the neutron star is scaled down strongly by the ratio $(\Omega_S/\Omega_{K*})^{4/3} \ll 1$ when the source resides on the propeller line and accretion is minimal.

The value of $\xi = 0.5$ used so far is justified by numerical simulations of star-disk models in low-power states near the propeller line ([Romanova et al. 2002, 2003, 2004](#); [Long et al. 2005](#); [Bessolaz et al. 2008](#); [Zanni & Ferreira 2013](#)). These simulations were designed to investigate protostellar systems and low-luminosity accreting neutron stars and they do not address the question of propeller emission at the Eddington limit ($L_X \approx L_{\text{Edd}}$). For such (ULX-type) systems in which minimal accretion proceeds at the Eddington rate, a value of $\xi = 1$ appears to be appropriate. This value is supported by a variety of theoretical calculations ([Wang 1987, 1995, 1996](#); [Arons 1993](#); [Ostriker & Shu 1995](#)). The stellar magnetic field on the propeller line of such systems is also expected to be stronger than that found on the lowest propeller line of Magellanic pulsars ([Christodoulou et al. 2016](#)), and the stronger magnetic field should not allow the disk to cross inside the spherical magnetospheric radius. For $\xi = 1$ and $\mu_* = \alpha B_* R_*^3$, eq. (16) takes a form independent of α , that is

$$L_X = \left(\frac{\eta}{\sqrt{2}}\right) B_*^2 R_*^3 \Omega_S \left(\frac{\Omega_S}{\Omega_{K*}}\right)^{4/3}, \quad (18)$$

and this equation allows for an independent determination of the stellar magnetic field B_* of ULX-type sources from measurements of Ω_S and the minimum L_X *without assuming torque balance*. Only one such ULX source is currently known to harbor a pulsar, M82 X-2, and we provide our estimates for this system in the next subsection.

2.6 Application to M82 X-2

The following measurements have been collected from observations of M82 X-2:

- (a) [Bachetti et al. \(2014\)](#) measured $P_S = 1.3725 \text{ s}$ and $\dot{P}_S = -2 \times 10^{-10} \text{ s s}^{-1}$ from *NuSTAR* observations when the system was in outburst.
- (b) [Brightman et al. \(2016\)](#) analyzed 15 years of *Chandra* observations and their results show that the source was sitting in its lowest-power state with $L_X = 1.0 \times 10^{38} \text{ erg s}^{-1}$ in at least two epochs (in 1999 and 2009), and perhaps in February 2013 too.
- (c) [Tsygankov et al. \(2016\)](#) analyzed the same *Chandra* observations and their results are similar reporting a lowest-power state with $L_X = (1.7 - 2.3) \times 10^{38} \text{ erg s}^{-1}$.

- (d) [Dall’Osso et al. \(2016\)](#) found an upper limit of $L_X = 1.7 \times 10^{38} \text{ erg s}^{-1}$ for no detection of the pulsar in a *Chandra* observation.

We assume that the lowest-power state found in the *Chandra* observations occurs when the source lies on its propeller line and we proceed to calculate its minimum X-ray luminosity and the stellar magnetic field. We use $G = 6.67 \times 10^{-8} \text{ cm}^3 \text{ g}^{-1} \text{ s}^{-2}$, $\xi = 1$, $I_* = 2MR_*^2/5$, and the following canonical pulsar parameters: $M = 1.4M_\odot$ and $R_* = 10 \text{ km}$.

Using the measured values of P_S and \dot{P}_S , we find from eq. (9) that

$$L_X = 7.1 \times 10^{38} \eta \text{ erg s}^{-1}. \quad (19)$$

The results of [Körding et al. \(2006\)](#) suggest that $\eta = 1/4$ is typical for neutron stars. Then, the power output of M82 X-2 on its propeller line appears to be equal to the Eddington limit.

Using the measured value of \dot{P}_S (i.e., torque balance), we also find from eq. (12) that

$$B_* = 3.1 \times 10^{12} \text{ G}. \quad (20)$$

But the magnetic field can also be determined independently from the minimum luminosity on the propeller line and without using the value of \dot{P}_S . In this case, eq. (18) can be written in the form

$$B_* = 8.0 \times 10^{11} \sqrt{\frac{L_X/(10^{38} \text{ erg s}^{-1})}{\eta}} \left(\frac{P_S}{1 \text{ s}}\right)^{7/6} \text{ G}. \quad (21)$$

For $P_S = 1.3725 \text{ s}$, $L_X = 1.0 \times 10^{38} \text{ erg s}^{-1}$ (based on the results of [Brightman et al. 2016](#)), and $\eta = 1/4$ this method yields

$$B_* = 2.3 \times 10^{12} \text{ G}. \quad (22)$$

Alternatively, for $L_X = L_{\text{Edd}} = 1.8 \times 10^{38} \text{ erg s}^{-1}$ (based on the results of [Tsygankov et al. 2016](#)), the stellar magnetic field is given again by eq. (20). The agreement between the above determinations of B_* justifies the implementation of torque balance in the derivation of eq. (20).

[Tsygankov et al. \(2016\)](#) found a bimodal distribution of X-ray luminosities in the *Chandra* data for M82 X-2 with a peak value of $L_X = 1.0 \times 10^{40} \text{ erg s}^{-1}$ in the high-power state and $L_X = 2.8 \times 10^{38} \text{ erg s}^{-1}$ in the low-power state. They assumed that the source straddles the [Corbet \(1996\)](#) gap between the two states and exhibits magnetospheric emission in the low-power state. We do not believe that this assessment of the two states of M82 X-2 is correct. For $P_S = 1.3725 \text{ s}$, the [Corbet \(1996\)](#) gap is given by

$$\Gamma = 167.87 P_S^{2/3} = 207, \quad (23)$$

and the low state below the gap is located at a luminosity of $L_X/\Gamma = 4.8 \times 10^{37} \text{ erg s}^{-1}$, if the propeller line is at the high observed state ($L_X = 1 \times 10^{40} \text{ erg s}^{-1}$). Then the low observed state lands well within the gap. Conversely, if the low observed state ($L_X = 2.8 \times 10^{38} \text{ erg s}^{-1}$) is assumed to be at the low boundary of the gap, then the propeller line is located 207 times higher at $L_X = 5.8 \times 10^{40} \text{ erg s}^{-1}$, i.e., several times higher than the most luminous observations to date ([Brightman et al. 2016](#); [Tsygankov et al. 2016](#)). In this case, the outbursts of the source land well within the

gap. For these reasons, our assumption that the low observed state corresponds to the propeller line and the high observed state corresponds to type-II outbursts (Coe et al. 2010) appears to be more reasonable.

3 DISCUSSION AND CONCLUSIONS

We have used standard accretion theory to estimate the typical values of the minimum isotropic luminosity and the surface magnetic field for the ULX pulsar M82 X-2 for which Bachetti et al. (2014) have recently measured pulsations with period $P_S = 1.3725$ s and spinup rate $\dot{P}_S = -2 \times 10^{-10}$ s s $^{-1}$. For these values of P_S and \dot{P}_S , our results show that the minimum luminosity due to accretion (the propeller line) is equal to the Eddington luminosity. We suspect that this also will turn out to be a characteristic property of other ULX sources that harbor pulsars.

X-ray emission at the Eddington level has been previously observed from A0538-66 in the Large Magellanic Cloud (Skinner et al. 1982) and the spin period of the pulsar (69.2 ms) was then measured. All other X-ray observations of A0538-66 since then, summarized by Kretschmar et al. (2004), have measured no pulsations and much smaller luminosities that fall well below the propeller line (Stella et al. 1986; Christodoulou et al. 2016). This led Campana et al. (1995) and Campana (1997) to propose that the low-power X-rays are emitted from the magnetosphere where accretion is halted by the centrifugal force and a luminous shock forms. M82 X-2, for which Kong et al. (2007) and Dall’Osso et al. (2016) have reported three cases of no detection, may also be caught in such nonpulsating sub-Eddington states if it will be observed below its propeller line ($L_X = 1 \times 10^{38}$ erg s $^{-1}$), and spectral differences (e.g., softer spectra) ought to provide an indication that the propeller line has been crossed. In fact, another ULX source hosting a stellar-mass black hole (NGC7793 P13) has been observed in such a “faint” state ($L_X = 5 \times 10^{37}$ erg s $^{-1}$; Motch et al. 2014).

The surface magnetic field of M82 X-2 is estimated to be $B_* \approx 3$ TG by two independent methods. One estimate comes from the lowest observed states (Brightman et al. 2016; Tsygankov et al. 2016) that effectively define the propeller line. This method does not use the measured value of the spinup rate \dot{P}_S , but it depends on the efficiency η of converting accretion power to X-rays (eq. [21]). The second estimate uses torque balance and $|\dot{P}_S|$, but it does not depend on η (eq. [20]). The results converge to effectively the same moderate value for B_* despite the uncertainties in the control parameters (ξ and η) and the direct measurements (P_S and \dot{P}_S) involved.

From the above estimates, we conclude that there is no need to invoke powerful, magnetar-type magnetic fields to explain the pulsar’s X-ray power output on the propeller line, standard accretion mechanisms suffice. Ekşi et al. (2015) and Lyutikov (2014) advocated for such very strong magnetic fields ($B \sim 10 - 100$ TG) because they thought that M82 X-2 is in spin equilibrium characterized by a substantial reduction in accretion torque. Our results in §§ 2.3–2.6 and the measured spinup rate \dot{P}_S argue against this claim and against using the peak X-ray luminosity in order to estimate the stellar magnetic field. This is also a

problem in the calculation of Tong (2015), although L_X was arbitrarily reduced from its peak value by an anisotropic factor of 5, causing a reduction to the dipolar B_* by 5^3 and making it agree with our determination (see also Pan et al. 2016, where the dipolar field undergoes decay during supercritical accretion and ends up at about the same value).

The studies by Kluźniak & Lasota (2015) and Dall’Osso et al. (2015, 2016) have assumed that the enormous power output of M82 X-1 is isotropic. We do not support these scenarios because, if that were the case, the radiation pressure could easily disrupt the accretion process and the inner disk itself. This point is made clear by the calculations of Pringle & Rees (1972) who found limits to the mass inflow rate in order for quasistatic accretion to continue uninterrupted by the radiation pressure. In the same context, we also do not agree with the argument of Ekşi et al. (2015) and Dall’Osso et al. (2015) that a very strong magnetic field would suppress the Thomson and Compton scattering cross-sections near the star’s surface permitting thus accretion to release energy above the conventional Eddington limit. We disagree because the magnetic field decreases quite rapidly with distance from the star and the problem of super-Eddington X-ray emission would then become acute at a few tens of stellar radii above the star’s surface where the magnetic field becomes weak and the scattering cross-sections assume again their standard values.

Thus it appears that a rather large (but theoretically attainable) amount of anisotropy appears to be necessary in order to explain the super-Eddington luminosities ($L_{\max} \approx 2 \times 10^{40}$ erg s $^{-1}$) observed during outbursts (Brightman et al. 2016; Tsygankov et al. 2016). In these episodic events, the supply of matter (the \dot{M}) at corotation increases by a factor of ~ 2 (Bachetti et al. 2014) and accretion becomes supercritical. The radiation cannot escape in all directions because the surrounding inflowing stream is optically thick. The photons can only escape in the axial direction where radiation pressure pushes an outflow to speeds larger than the local escape speed. The outflow is highly anisotropic and we need to be in a favored position in order to detect the apparent super-Eddington flux liberated in the funnel. If we assume that a luminosity of L_{Edd} appears to the observer as L_{\max} , then we find that the opening angle of the emission funnel pointing in our direction is about 6.5° . The fact that we need to be in such a favored position explains the scarcity of ULX sources.

Furthermore, if the emission from M82 X-2 is collimated, some of the energy must emerge at much longer wavelengths. This is the case according to the radio observations of Kronberg et al. (1985), McDonald et al. (2002), and Fenech et al. (2008); and the infrared observations of Kong et al. (2007) and Gandhi et al. (2011). The radio maps show a core-dominated source which is expected if the pulsar is a modestly aligned rotator and a collimated jet is coming out in our direction. Kong et al. (2007) also produced *Chandra* X-ray spectra that are hard (photon indices 1.3 – 1.7 from an absorbed power-law model) and show no soft excess. This, combined with the strong X-ray variability on timescales of ~ 2 months and the recurring type-II outbursts, indicates that M82 X-2 is not at all dissimilar to Galactic and Magellanic X-ray binaries harboring neutron stars. The recent observations reported by

Brightman et al. (2016) and Tsygankov et al. (2016) have effectively confirmed this picture.

There are some more discoveries about ULX sources indicating that they may be supercritical accretors and type-II bursters, but otherwise not at all exotic:

(a) Gilfanov et al. (2004) found that the ULX sources are just the high end of a luminosity function that cuts off at $L_X \sim 3 \times 10^{40} \text{ erg s}^{-1}$ and in which the known high-mass X-ray binaries make up the low end of a single power-law with slope ~ 1.6 .

(b) Luangtip et al. (2014) found that the spectral index of ULX spectra changes around $L_X \sim 2 \times 10^{39} \text{ erg s}^{-1}$, a value that may indicate a transition from critical to supercritical accretion by common $10M_\odot$ black holes.

(c) Laycock et al. (2015) determined from the radial velocity curve of IC10 X-1 that the compact object is likely a neutron star, although a low-stellar-mass black hole cannot be ruled out.

(d) Liu et al. (2013) determined from optical spectroscopy that the mass of the compact object in M101 ULX-1 is no more than $30M_\odot$. As in the case of NGC7793 P13 (Motch et al. 2014), it is unlikely that this is an intermediate-mass black hole. For this system, Shen et al. (2016) developed a model that explains how reprocessed soft X-rays can be finally emitted from very large radii (~ 100 times beyond the inner radius of the accretion disk). This model may also be applicable to other ULX sources harboring stellar compact objects (see also King 2008).

Soon after the discovery of pulsations in M82 X-2, Doroshenko et al. (2015) tried to find more pulsating ULX sources in archival *XMM-Newton* observations, but they did not succeed. We believe that it is a matter of time until another pulsating ULX source is found and IC10 X-1 is an intriguing candidate. But because such powerful X-ray emission ($L_X > 10^{40} \text{ erg s}^{-1}$) from neutron stars must occur at a favorable angle to the observer, we expect very few such pulsars to be discovered in the future (see also King 2009).

Two very recent studies of M82 X-2 have produced results that are different by an order of magnitude and on opposite sides of our result ($B_* \approx 3 \text{ TG}$), but we understand the source of the differences. In both studies, the mass accretion rate \dot{M} at the disk's inner edge is too high. In King & Lasota (2016), the inner edge of the disk R_d is located at r_m which is 10 times smaller than corotation r_c . This causes highly supercritical accretion in the region between these two radii that can only occur if the magnetic field is weaker by a factor of 30. Thus, this model is not on the propeller line of M82 X-2.

In Karino & Miller (2016), the deduced $\dot{M} = 3.2 \times 10^{20} \text{ g s}^{-1}$ is even more uncertain because it depends on a host of assumptions about the unknown parameters of the donor star and the spherically symmetric mass transfer between the two stars at large radii; it turns out to be ~ 10 times larger than that determined in the King & Lasota (2016) model despite the fact that it was determined at corotation assuming that $r_c = r_m$. Thus, this model is not on the propeller line either. In both cases, the Eddington value of $\dot{M}_{\text{Edd}} = 2 \times 10^{18} \text{ g s}^{-1}$ applied at the disk's inner edge R_d produces our result. Although it is not obvious how such a value can be deduced from the Karino & Miller (2016) model (too many unknown parameters are involved), the

King & Lasota (2016) model, which relies on the anisotropy of the X-ray emission, can be brought to agreement if the location of R_d is moved closer to corotation.

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